

Single-Transverse Spin Asymmetry in Dijet Correlations at Hadron Colliders

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Abstract

We present a phenomenological study of the single-transverse spin asymmetry in azimuthal correlations of two jets produced nearly “back-to-back” in pp collisions at RHIC. We properly take into account the initial- and final-state interactions of partons that can generate this asymmetry in QCD hard-scattering. Using distribution functions fitted to the existing single-spin data, we make predictions for various weighted single-spin asymmetries in dijet correlations that are now readily testable at RHIC.

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I. INTRODUCTION

Single-transverse spin asymmetries (SSAs) play a fundamental role for our understanding of QCD in high-energy hadronic scattering. They may be obtained for reactions in, for example, lepton-proton or proton-proton scattering with one transversely polarized initial proton, by dividing the difference of the cross sections for the two settings of the transverse polarization by their sum. There have been extensive experimental investigations of such asymmetries [1, 2, 3, 4, 5, 6, 7]. These have initiated much theoretical progress, in particular within the last few years.

A particular focus has been on a class of single-spin observables that are characterized by a large momentum scale Q (for example, the virtuality of the photon in deeply-inelastic scattering (DIS)) and by a much smaller, but also measured, transverse momentum q_\perp . In such a “two-scale” situation, single-spin asymmetries may arise at leading power, that is, not suppressed by an inverse power of Q . For some of these cases, factorization theorems have been established [8, 9, 10] that allow to write the spin-dependent cross sections in terms of parton distribution functions and/or fragmentation functions, perturbative hard-scattering functions, and so-called soft factors. A crucial feature is that the distribution functions and the soft factor in this factorization are not integrated over the transverse momenta of partons, because these in fact generate the observed transverse momentum q_\perp . Among other things, the observables may therefore provide valuable insights into the dependence of parton distributions in nucleons on transverse momentum. This becomes particularly interesting when the nucleon is transversely polarized, because there may be correlations between the nucleon spin vector, its momentum, and the parton’s transverse momentum. One particular correlation, known as the “Sivers effect” and described by so-called “Sivers functions” [11], is now widely believed to be involved in a variety of observed hadronic single-spin phenomena.

Closer theoretical studies have revealed that the Sivers effect plays an important role in QCD, beyond giving rise to phenomenological functions to be used in the description of single-spin asymmetries. A particularly interesting feature is that the Sivers effect is not universal in the usual sense, that is, it is not represented by universal probability functions convoluted with partonic hard-scattering cross sections. This might at first sight appear to make the study of these functions less interesting. However, the non-universality has in fact a clear physical origin, and its closer investigation has turned out to be an extremely

important and productive development in QCD. In a nutshell, in order not to be forced to vanish because of the time-reversal symmetry of QCD, single-spin asymmetries require the presence of a strong-interaction phase. For the Sivers functions this phase originates from the “gauge links” in their definition [12, 13, 14], which are path-ordered exponentials of the gluon field that make the functions gauge-invariant. In DIS, the gauge link may be viewed as a rescattering of the parton in the color field of the nucleon remnant. That such a final-state rescattering may generate the phase required for a nonzero SSA in semi-inclusive hadron production in DIS (SIDIS) was first discovered within a model calculation [15].

Depending on the hard-scattering process, the “rescattering” will however manifest itself in different ways. For example, for Drell-Yan lepton pair production in hadronic scattering, *initial-state*, rather than final-state, interactions are relevant. As a result, the phase provided by the gauge links is opposite, and the Sivers functions for the Drell-Yan process have opposite sign [12, 13, 14, 15]. In more general terms, the nontrivial “universality” property of the Sivers functions is the direct consequence of gauge interactions in Quantum Chromodynamics [12, 13, 14], and of the QCD factorization theorems applying to the relevant hard processes [8, 9, 10]. It is a remarkable and fundamental QCD prediction that really tests all concepts we know of for analyzing hard-scattering reactions in strong interactions, and it awaits experimental verification.

While measurements of SSAs in SIDIS are now maturing and have established the presence of Sivers-type contributions [6], it will still be a while until precise single-spin measurements in the relatively rare Drell-Yan process will become feasible at RHIC [16] or elsewhere [17]. However, there are of course other hard-scattering reactions in hadronic collisions for which single-transverse spin asymmetries may be defined, and that may potentially be used in lieu of the Drell-Yan process for testing the nontrivial “universality” properties of the Sivers functions. In [18], it was proposed to use the SSA in azimuthal correlations of two jets produced in pp collisions at RHIC to learn about the Sivers functions. To a first approximation, such dijets are produced by $2 \rightarrow 2$ partonic QCD hard-scattering. With collinear kinematics, the jets are exactly “back-to-back” in the plane perpendicular to the initial beam directions and thus separated by 180° in azimuthal angle in this plane. Partonic transverse momenta will generate deviations from this topology, because they will lead on average to a non-vanishing net transverse momentum q_\perp of the jet pair, much smaller than each of the jet transverse momenta P_\perp individually. The observable is therefore of the

“two-scale” type described above, and as was shown in [18], if one proton is polarized, a single-spin asymmetry may be defined that is in principle sensitive to the Sivers functions. As a caveat, factorization of the spin-dependent cross section for this observable in terms of transverse-momentum-dependent (TMD) functions still remains to be established.

Unlike the relatively simple cases of SIDIS and the Drell-Yan process, where either final-state or initial-state interactions contribute to the Sivers asymmetry, both are present for dijet production [19]. This complicates the analysis of the process-dependence of the Sivers functions considerably, but at the same time it also makes it much more interesting from a theoretical point of view, because the interactions in this case are “truly QCD”, that is, they involve the detailed gauge structure of the theory, including for example its non-abelian nature. At the time of [18], the process-dependence had not yet been worked out for the case of the SSA in dijet production, so that phenomenological predictions had to remain qualitative, at best. Over the last year, however, there has been extensive work on deriving and clarifying the structure of the gauge links for this and related processes in pp collisions [20, 21, 22]. Indeed, the resulting structure is far more complicated than that in SIDIS or the Drell-Yan process. However, it turns out that if one takes a certain weighted integral (“moment”) of the asymmetry, remarkable simplifications occur. This moment is defined by integrating the spin-dependent cross section with a factor $\sin \delta$, where δ is the azimuthal imbalance between the two jets ($\delta = 0$ if the jets are exactly back-to-back in azimuth). For each of the various contributing $2 \rightarrow 2$ partonic channels, the gauge link then essentially collapses into a set of simple color factors that multiply contributions from color-gauge invariant subamplitudes to the given partonic process. One may, in fact, for convenience choose to absorb these factors into the hard-scattering functions, and *define* the Sivers functions as the functions measured in the SSA in SIDIS. In this way, the net effect of the gauge links on the $\sin \delta$ -moment of the spin-dependent cross section is to generate new partonic hard-scattering functions that are different from the usual spin-independent ones, but that are actually similarly simple in structure.

At the same time, taking the moment of the factorized spin-dependent cross section leads to a new expression that involves only a certain moment of the Sivers functions in partonic transverse momentum k_\perp , rather than the fully transverse-momentum dependent functions themselves. As was shown in [14], these k_\perp -moments of the Sivers functions are directly related to twist-three quark-gluon correlation functions first introduced in [23, 24] to

describe the SSA for single-inclusive hadron production in hadronic scattering. By now, quite some knowledge about these correlation functions has been gathered from phenomenological studies [25] of the corresponding data.

The upshot of all this is that it has now become possible for the first time to make predictions for the $\sin \delta$ -moment of the single-transverse spin asymmetry in dijet production at RHIC that correctly take into account the process-dependence of the Sivers functions and incorporate phenomenological information on some properties of the functions that is available from other measurements. This is the goal of this note.

II. SPIN-DEPENDENT CROSS SECTION AND $\sin \delta$ -MOMENT

We study azimuthal correlations of two jets produced nearly “back-to-back” in a hadronic collision. More specifically, we are interested in situations in which the sum of the two jet transverse momenta, $\vec{q}_\perp \equiv \vec{P}_{1\perp} + \vec{P}_{2\perp}$ (or a component or projection thereof), is measured and small, while both individual jet transverse momenta are large and similar. We will therefore approximate $\vec{P}_{1\perp} = -\vec{P}_{2\perp} \equiv \vec{P}_\perp$ wherever possible. For the lengths of these momentum vectors we will simply write $P_\perp = |\vec{P}_\perp|$ and $q_\perp = |\vec{q}_\perp|$. The differential cross section for the process with one transversely polarized hadron contains terms of the form

$$\frac{2\pi d^6\sigma}{d\eta_1 d\eta_2 dP_\perp^2 d\phi_1 d^2q_\perp} = \frac{2\pi d^6\sigma_{UU}}{d\eta_1 d\eta_2 dP_\perp^2 d\phi_1 d^2q_\perp} + \hat{e}_z \cdot (\vec{S}_\perp \times \hat{q}_\perp) \frac{2\pi d^6\sigma_{TU}}{d\eta_1 d\eta_2 dP_\perp^2 d\phi_1 d^2q_\perp}, \quad (1)$$

where \hat{e}_z is a unit vector in the direction of the polarized proton, \vec{S}_\perp is the transverse spin vector of the polarized proton, and η_1 and η_2 are the pseudo-rapidities of the two jets. The first term in Eq. (1) represents the unpolarized (or spin-averaged) cross section, while the second term is the single-transverse-spin dependent one. We note that the angular dependence of the spin-dependent term is $|\vec{S}_\perp| \sin(\phi_b - \phi_S)$, where $\phi_b = (\phi_1 + \phi_2)/2$ is the so-called bisector-angle of the two jets, which (approximately) corresponds to the direction of \vec{q}_\perp . For this reason one may also choose to integrate the six-fold differential cross section in Eq. (1) over ϕ_1 , keeping \vec{q}_\perp fixed.

As a generalization of the SIDIS and Drell-Yan cases [9], we can write down a factorization formula for the differential cross section at small imbalance ($\vec{q}_\perp \neq 0$) of the jets, in terms of TMD parton distributions, soft factors, and hard-scattering functions [26]. We remind the reader that such a factorization still remains to be proven. In this paper, we will not

discuss the details of factorization issues related to the dijet correlations. As we mentioned above, significant simplifications occur when the imbalance of the two jets is integrated out by taking certain moments. For example, integrating the spin-independent differential cross section over all \vec{q}_\perp , its expression reverts to the standard collinear factorization formula for dijet production,

$$\begin{aligned} \langle 1 \rangle_{UU} &\equiv \int d^2 \vec{q}_\perp \frac{2\pi d^6 \sigma_{UU}}{d\eta_1 d\eta_2 dP_\perp^2 d\phi_1 d^2 \vec{q}_\perp} = \frac{2\pi d^4 \sigma_{UU}}{d\eta_1 d\eta_2 dP_\perp^2 d\phi_1} \\ &= \sum_{ab} x_a f_a(x_a) x_b f_b(x_b) H_{ab \rightarrow cd}^{uu}(\hat{s}, \hat{t}, \hat{u}) , \end{aligned} \quad (2)$$

where the f_i denote the usual collinear (light-cone) parton distribution functions for parton type $i = q, \bar{q}, g$. We have assumed here that these are recovered by integration of the corresponding TMD distribution functions over all partonic transverse momentum, and we disregard the renormalization properties of the operators defining these distributions and the scale dependence introduced by renormalization. However, all these effects can be systematically included accordingly.

The factors $H_{ab \rightarrow cd}^{uu}$ in Eq. (2) are the customary hard-scattering cross sections $d\hat{\sigma}^{ab \rightarrow cd}/d\hat{t}$ for the partonic processes $ab \rightarrow cd$ (for a compilation, see, for example, Ref. [22]). They are functions of the partonic Mandelstam variables, $\hat{s} = (p_a + p_b)^2$, $\hat{t} = (p_a - p_c)^2$, $\hat{u} = (p_a - p_d)^2$, with obvious notation of the parton momenta. In terms of the hadronic center-of-mass energy \sqrt{s} and the jet transverse momenta and pseudo-rapidities, one has $\hat{s} = x_a x_b s$, $\hat{t} = -P_\perp^2 (e^{\eta_2 - \eta_1} + 1)$, $\hat{u} = -P_\perp^2 (e^{\eta_1 - \eta_2} + 1)$, where the partonic momentum fractions are fixed by $x_a = \frac{P_\perp}{\sqrt{s}} (e^{\eta_1} + e^{\eta_2})$, $x_b = \frac{P_\perp}{\sqrt{s}} (e^{-\eta_1} + e^{-\eta_2})$.

Next we turn to the single-transverse-spin dependent differential cross section $d\sigma_{TU}$ in Eq. (1), which can be further simplified by taking a moment in

$$P_\perp \sin \delta = \frac{\hat{e}_z \cdot (\vec{P}_\perp \times \vec{q}_\perp)}{P_\perp} , \quad (3)$$

where $\delta = \pi - (\phi_2 - \phi_1)$ measures how far the two jets are away from the back-to-back configuration. Within our approximations, the weight $P_\perp \sin \delta$ corresponds to a weight in q_\perp . One finds [21]:

$$\begin{aligned} \left\langle \frac{2 P_\perp \sin \delta}{M_P} \right\rangle_{TU} &\equiv |\vec{S}_\perp| \int d^2 \vec{q}_\perp \frac{q_\perp}{M_P} \frac{2\pi d^6 \sigma_{TU}}{d\eta_1 d\eta_2 dP_\perp^2 d\phi_1 d^2 \vec{q}_\perp} \\ &= |\vec{S}_\perp| \sum_{ab} x_a \frac{-1}{M_P} g T_{Fa}(x_a) x_b f_b(x_b) H_{ab \rightarrow cd}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u}) , \end{aligned} \quad (4)$$

where M_P is the proton mass, g is the strong coupling constant and the $T_{Fa}(x)$ are the Qiu-Sterman matrix elements or quark-gluon correlation functions [24], defined as

$$T_{Fa}(x) = \int \frac{d\xi^- d\eta^-}{4\pi} e^{i(xP^+ \eta^-)} \epsilon_{\perp}^{\beta\alpha} S_{\perp\beta} \langle PS | \bar{\psi}^a(0) \gamma^+ F_{\alpha}^+(\xi^-) \psi^a(\eta^-) | PS \rangle, \quad (5)$$

with the quark fields ψ^a (for flavor a) and the gluon field strength tensor F_{α}^+ . The $T_{Fa}(x)$ matrix elements enter because they are related to the k_{\perp} -moment of the (SIDIS) Siverson function for quark flavor a [14], that is, $g T_{Fa}(x) = -2M f_{1T,a}^{\perp(1)}(x)$. We note that there could also be contributions by purely gluonic “ ggg ” correlation functions. These will be ignored for now, so that the label a in Eq. (4) runs only over quarks and anti-quarks. Furthermore, there is actually a second contribution to the single-transverse-spin dependent cross section, which involves the scattering of transversely polarized quarks from both the polarized proton (transversity distribution, δf or h_1) and from the unpolarized proton (Boer-Mulders function h_1^{\perp} [27]). The latter functions are also effects of initial- and final-state interactions and appear in a matrix element for unpolarized protons similar to Eq. (5). Like the Siverson-type “ ggg ” correlation functions, we will also ignore the contributions associated with the Boer-Mulders functions in the present study, even though a future more detailed analysis of forthcoming experimental data may well require to take all of these into account.

The relevant hard-scattering functions $H_{ab \rightarrow cd}^{\text{sivers}}$ have been calculated in [21, 22], where they were termed “gluonic-pole cross sections” due to their association with the Qiu-Sterman or gluonic pole matrix elements [24]. We have also independently reproduced [28] the $H_{ab \rightarrow cd}^{\text{sivers}}$ within a model calculation along the lines of Ref. [15]. For convenience, we list the ones that will be relevant for our numerical calculations presented below:

$$\begin{aligned} H_{qq' \rightarrow qq'}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u}) &= \frac{\alpha_s^2 \pi}{\hat{s}^2} \frac{N_c^2 - 5}{4N_c^2} \frac{2(\hat{s}^2 + \hat{u}^2)}{\hat{t}^2}, \\ H_{q\bar{q}' \rightarrow q\bar{q}'}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u}) &= \frac{\alpha_s^2 \pi}{\hat{s}^2} \left(-\frac{N_c^2 - 3}{4N_c^2} \right) \frac{2(\hat{s}^2 + \hat{u}^2)}{\hat{t}^2}, \\ H_{q\bar{q} \rightarrow q'\bar{q}'}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u}) &= \frac{\alpha_s^2 \pi}{\hat{s}^2} \frac{N_c^2 + 1}{4N_c^2} \frac{2(\hat{t}^2 + \hat{u}^2)}{\hat{s}^2}, \\ H_{qq \rightarrow qq}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u}) &= \frac{\alpha_s^2 \pi}{\hat{s}^2} \left\{ \frac{N_c^2 - 5}{4N_c^2} \left[\frac{2(\hat{s}^2 + \hat{u}^2)}{\hat{t}^2} + \frac{2(\hat{s}^2 + \hat{t}^2)}{\hat{u}^2} \right] + \frac{N_c^2 + 3}{4N_c^3} \frac{4\hat{s}^2}{\hat{t}\hat{u}} \right\}, \\ H_{q\bar{q} \rightarrow q\bar{q}}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u}) &= \frac{\alpha_s^2 \pi}{\hat{s}^2} \left\{ -\frac{N_c^2 - 3}{4N_c^2} \frac{2(\hat{s}^2 + \hat{u}^2)}{\hat{t}^2} + \frac{N_c^2 + 1}{4N_c^2} \frac{2(\hat{u}^2 + \hat{t}^2)}{\hat{s}^2} - \frac{N_c^2 + 1}{4N_c^3} \frac{4\hat{u}^2}{\hat{s}\hat{t}} \right\}, \end{aligned}$$

$$\begin{aligned}
H_{qg \rightarrow qg}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u}) &= \frac{\alpha_s^2 \pi}{\hat{s}^2} \left\{ -\frac{N_c^2}{4(N_c^2 - 1)} \frac{2(\hat{s}^2 + \hat{u}^2)}{\hat{t}^2} \left[\frac{\hat{s}}{\hat{u}} - \frac{\hat{u}}{\hat{s}} \right] - \frac{1}{2(N_c^2 - 1)} \frac{2(\hat{s}^2 + \hat{u}^2)}{\hat{t}^2} \right. \\
&\quad \left. - \frac{1}{4N_c^2(N_c^2 - 1)} \frac{2(\hat{s}^2 + \hat{u}^2)}{\hat{s}\hat{u}} \right\}, \\
H_{q\bar{q} \rightarrow gg}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u}) &= \frac{\alpha_s^2 \pi}{\hat{s}^2} \left\{ -\frac{1}{2N_c} \frac{2(\hat{t}^2 + \hat{u}^2)}{\hat{s}^2} + \frac{N_c}{4} \frac{2(\hat{t}^2 + \hat{u}^2)}{\hat{s}^2} \left[\frac{\hat{t}}{\hat{u}} + \frac{\hat{u}}{\hat{t}} \right] \right. \\
&\quad \left. - \frac{2N_c^2 + 1}{4N_c^3} \frac{2(\hat{t}^2 + \hat{u}^2)}{\hat{t}\hat{u}} \right\}, \tag{6}
\end{aligned}$$

where $N_c = 3$ is the number of colors. Hard-scattering functions corresponding to gluonic correlation functions were also calculated in [22], but are not taken into account in our present study as we discussed above. Comparing the above functions with the usual spin-averaged hard-scattering functions [22], one can see that they essentially differ in the color factors that multiply terms with similar kinematic structure. This is the net result of the combined effects from the initial and final state interactions in dijet production in hadronic reactions.

III. PHENOMENOLOGICAL RESULTS FOR RHIC

We now use the above formulas to obtain some predictions for the SSA in dijet-production at RHIC. A SSA for our $(2P_\perp \sin \delta / M_P)$ -moment (i.e. q_\perp / M_P -moment) can be defined as (from now on we choose $|\vec{S}_\perp| = 1$)

$$\frac{\langle 2P_\perp \sin \delta / M_P \rangle_{TU}}{\langle 1 \rangle_{UU}} = \frac{\sum_{ab} x_a \frac{-1}{M_P} g T_{Fa}(x_a) x_b f_b(x_b) H_{ab \rightarrow cd}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u})}{\sum_{ab} x_a f_a(x_a) x_b f_b(x_b) H_{ab \rightarrow cd}^{uu}(\hat{s}, \hat{t}, \hat{u})}. \tag{7}$$

The quark-gluon correlation functions T_{Fa} have recently been fitted [25] to data [1, 2, 5] for the SSA in single-inclusive hadron production in hadronic collisions. Two sets of T_{Fa} were presented in [25]. For definiteness, we choose set I, which is a two-flavor fit with valence u and d -quark T_F distributions only. For these the following parameterizations were obtained in [25]:

$$T_{Fu_v}(x) = 0.275x^{0.508}(1-x)^{0.399}u_v(x), \quad T_{Fd_v}(x) = -0.365x^{-0.108}(1-x)^{0.287}d_v(x), \tag{8}$$

where the u_v, d_v are the corresponding unpolarized valence quark distributions. For the latter, as for all other unpolarized parton distributions we need, we choose the CTEQ5L parameterizations [29]. We plot the resulting weighted asymmetry in Fig. 1. The kinematics are chosen to correspond to current measurements at STAR: both jets have transverse

momenta $5 \text{ GeV} \leq P_\perp \leq 10 \text{ GeV}$ and pseudo-rapidities $-1 \leq \eta_i \leq 2$. The asymmetry is plotted as a function of the sum of the two jet pseudo-rapidities. We have chosen the hard scale $\mu = P_\perp$ in the strong coupling constant and the unpolarized parton distributions.

For comparison, we also show in Fig. 1 the asymmetries that one would have expected if the Sivers functions relevant for dijet production were identical to the ones in SIDIS or the Drell-Yan process. In the framework of our present calculation, the corresponding partonic hard-scattering functions would in this case be identical to the spin-averaged functions $H_{ab \rightarrow cd}^{uu}$, or to their negatives. The dotted-dashed lines in the right panel of Fig. 1 represent these cases. They essentially correspond to the earlier predictions of [18] and [26] that were made when the process-dependence of the Sivers functions had not yet been worked out for the dijet case. One can see that when the correct process-dependence is incorporated, the asymmetry overall becomes smaller by about a factor two, which can be traced back to the color factors for the dominant subprocess $qg \rightarrow qg$. The sign is identical to the case when the Sivers functions for dijet production are assumed to be “DIS-like”, implying that in a sense final-state interactions dominate over initial-state ones.

In principle, one might verify experimentally the process-dependence of the Sivers functions by discriminating between the various curves in Fig. 1 and confirming the QCD-predicted result shown by the solid line. In practice, this may require good knowledge of the T_F distributions, and an understanding of issues like scale evolution and higher-order corrections. A closer analysis reveals that the asymmetry in Fig. 1 is the result of rather significant cancellations between contributions of opposite signs by u and d quarks. To show this, we display their individual contributions separately in the left panel of Fig. 1.

We finally also briefly discuss a related type of SSA in dijet production at RHIC. In [26], a differently weighted SSA was considered, defined in the following way:

$$\langle \text{Sgn}(\delta) \rangle_{TU} \equiv \int d^2 \vec{q}_\perp \text{Sgn}(\delta) \frac{d^5 \sigma_{TU}}{d\eta_1 d\eta_2 dP_\perp^2 d^2 \vec{q}_\perp}, \quad (9)$$

where $\text{Sgn}(x)$ is the sign function. This asymmetry has the property that the weight only depends on the azimuthal separation of the two jets, but not on their transverse momenta. This may be an advantage for experimental measurements when the jet energy scale is not precisely known. When applied to the TMD factorized expression for the spin-dependent cross section, the moment defined in Eq. (9) leads to an expression different from (4). One finds in fact that in general the transverse-momentum dependences of the various functions

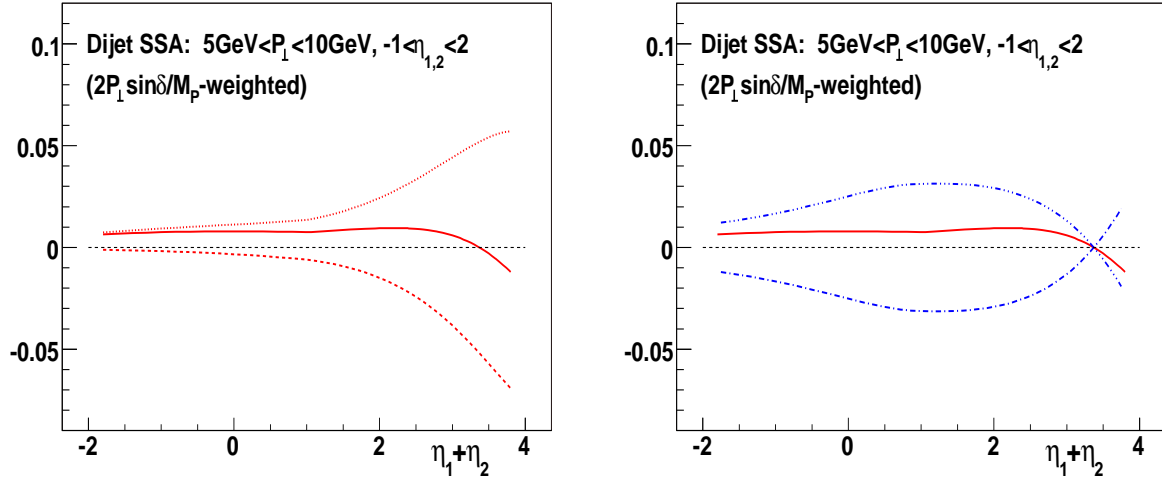


FIG. 1: The weighted single-transverse-spin asymmetry for dijet-correlations in polarized proton-proton scattering at RHIC, as a function of the sum of the jet pseudo-rapidities. The jet transverse momenta are in the range 5 – 10 GeV. The solid line is our main result, based on Eqs. (6)-(8). In the left panel we also plot the individual contributions due to the u-quark (dashed line) and d-quark (dotted line) Qiu-Sterman matrix elements. For comparison, in the right panel we also show the results that would be obtained if the Siverson functions contributing to dijet production were the same as those in SIDIS (upper curve) or the Drell-Yan process (lower curve).

do not completely decouple anymore, so that the final result can in general no longer be expressed in terms of only functions of light-cone momentum fractions. If one assumes for simplicity [26], however, that the only relevant dependence on transverse momentum resides in the Siverson functions, the resulting expression again resembles a collinearly-factorized one:

$$\frac{\langle \text{Sgn}(\delta) \rangle_{TV}}{\langle 1 \rangle_{UU}} = \frac{\sum_{ab} x_a q_{Ta}^{(1/2)}(x_a) x_b f_b(x_b) H_{ab \rightarrow cd}^{\text{sivers}}(\hat{s}, \hat{t}, \hat{u})}{\sum_{ab} x_a f_a(x_a) x_b f_b(x_b) H_{ab \rightarrow cd}^{uu}(\hat{s}, \hat{t}, \hat{u})}, \quad (10)$$

where the $q_{Ta}^{(1/2)}$ are the so-called “1/2-moments” of the Siverson functions for SIDIS and are defined as

$$q_{Ta}^{(1/2)}(x) \equiv \int d^2 k_{\perp} \frac{|\vec{k}_{\perp}|}{M} q_{Ta}^{\text{SIDIS}}(x, k_{\perp}). \quad (11)$$

In the above equations, we have used the notation “ q_{Ta} ” for the Siverson function for quark flavor a ; its definition is identical to that in [14]: $q_{Ta} \equiv f_{1T,a}^{\perp}$. In [26], the hard-scattering functions were chosen to be the same as the spin-averaged ones, with opposite sign. In light of the above discussions, one now would like to update the predictions for the asymmetry

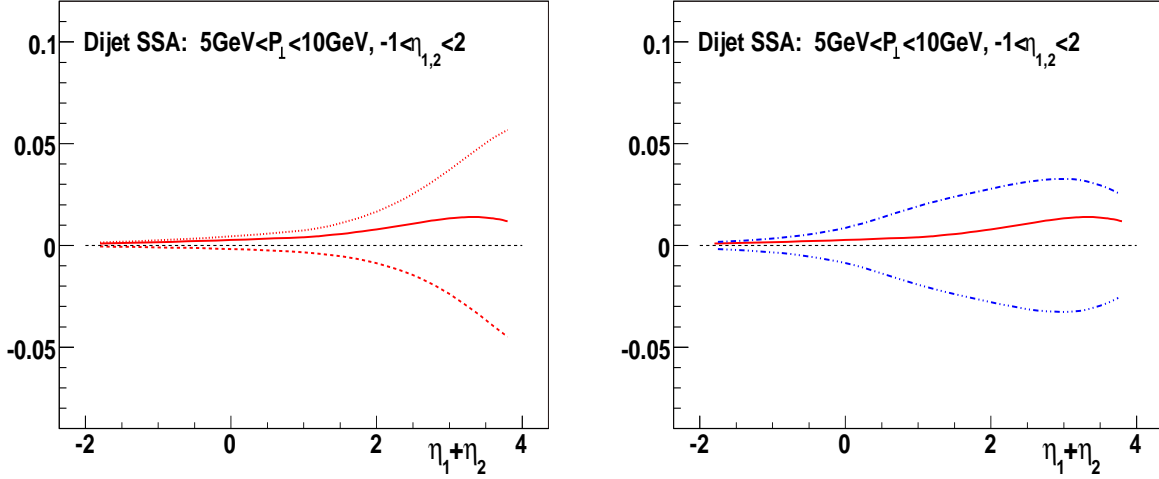


FIG. 2: Same as Fig. 1, but for the SSA defined in Eq. (10).

in (10). We will use the $H_{ab \rightarrow cd}^{\text{sivers}}$ given in Eqs. (6). We note that it remains to be established that the same $H_{ab \rightarrow cd}^{\text{sivers}}$ do contribute in this case. This is not a priori clear, because the general gauge link structure is very complex in the general TMD case, and it has so far only been demonstrated that the $H_{ab \rightarrow cd}^{\text{sivers}}$ apply when the $\sin \delta$ moment is taken. For now, we just conjecture that use of the $H_{ab \rightarrow cd}^{\text{sivers}}$ of [21, 22] is justified in this case; a closer discussion of this issue is left for future work [28].

The 1/2-moments of the quark Siverson functions were determined in [26] by a fit to the HERMES data [6] for the Siverson-type single-spin asymmetry:

$$u_T^{(1/2)}(x) = -0.75x(1-x)u(x), \quad d_T^{(1/2)}(x) = 2.76x(1-x)d(x). \quad (12)$$

These results correspond to set II presented in [26]. In Fig. 2, we show predictions for the asymmetry for the dijet-correlation defined in (10) at RHIC, based on this parameterization. We find that the asymmetry shares many features with the one shown for the $(2P_\perp \sin \delta / M_P)$ -moment in Fig. 1. We note that first preliminary experimental data for this asymmetry have now been presented by STAR [3], which are so far consistent with a vanishing asymmetry.

IV. CONCLUSIONS

In summary, we have studied single-transverse spin asymmetries in dijet-correlations at RHIC, making use of the recently derived partonic hard-scattering cross sections that properly incorporate the initial- and final-state interactions, and of distribution functions

fitted to existing data for single-spin asymmetries. We have found that the initial- and final-state interactions tend to decrease the magnitude of the asymmetry at RHIC with respect to earlier estimates that assumed the Sivers functions for this observable to be identical to the Sivers functions in SIDIS or the Drell-Yan process. Overall, the resulting asymmetries turn out to be more dominated by the final-state interactions, and hence more “DIS-like”.

With experimental data on dijet single-spin asymmetries now forthcoming, it will be interesting to perform detailed comparisons with the theoretical expectations. Other observables, such as the SSAs in dihadron production $pp \rightarrow h_1 h_2 X$ or in photon-plus-jet production $pp \rightarrow \gamma \text{jet } X$, will also be extremely interesting. It will be important to further develop the theoretical framework for all these observables, by addressing issues like TMD factorization, higher orders, soft factors, and Sudakov suppression, in particular.

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